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A linear theory of bunching in the orotron is derived in which the electron beam evolves in time via the Boltzmann equation. The theory is used to derive the power radiated by an electron beam subjected to a general longitudinal electric field. It is then specialized to the case of the orotron, in which the electric field is quasi-periodic. The starting current, I_S, and the electronic tuning transconductance, df/dV, are calculated

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20. Abstract (Cont'd)

from the theory. These results are evaluated for the planned Harry Diamond Laboratories orotron experiment at 75 gigahertz and are found to be I_s = 34.5 milliamperes and df/dV = 0.31 megahertz/volt.

CONTENTS

FIGURES

1. INTRODUCTION

The orotron[1-3] is a new type of electron device for generation and amplification of millimeter wave radiation. The device consists of an electron beam generator and collector and a Fabry-Perot resonator containing one grooved mirror (grating) and one smooth mirror. The principle of operation of the orotron is based on the Smith-Purcell effect.[4]

A previous study* described the physical principles of orotron operation and surveyed the existing experimental and theoretical literature (mostly from the Soviet Union). To augment the information in that report, quantify its conclusions, and aid in the development of the first operating orotron in the United States (at the Harry Diamond Laboratories — HDL), this series of reports will cover theoretical aspects of the orotron problem. A previous report[5] considered an equivalent circuit of the orotron and derived the output power and frequency characteristics of the device. This second report deals with the microscopic mechanism of electron beam bunching in the device.

Earlier bunching theories[6,7] considered an idealization of the electromagnetic field distribution near the diffraction grating. This approximation leads to an order-of-magnitude estimate of the threshold current, but does not give the correct dependence on electron beam thickness and grating parameters. In addition, the early theories considered the motion of a single electron according to Newton's second law. Although this approach is correct, it is unwieldy and does not lend itself to straightforward generalization to a nonlinear theory.

We consider the bunching problem from the point of view of distribution functions that describe the electron beam as a whole, and we consider the time evolution of the beam via the Boltzmann equation. This method provides a connection between the orotron and other free-electron lasers that have been analyzed by this method.[8,9] In section 2, a general formula is derived for bunching of an electron beam by an applied longitudinal rf electric field. Section 3 applies this result to the orotron and computes the output power of the device; by equating the power radiated by the electron beam to the power absorbed by the open resonator, we derive an expression for the starting current for the device. Section 4 justifies the representation of the electron beam as a complex frequency-dependent admittance (in the language of our earlier report[5]) and derives the electronic tuning transconductance of the orotron.

2. ELECTRON MOTION IN LONGITUDINAL ELECTRIC FIELD

In the orotron (as opposed to other free-electron laser devices),[8,9] electrons are bunched by the longitudinal rf electric field (the component of the electric field along the direction of electron motion) in the open resonator. In this section, we analyze the motion of an electron beam in a monochromatic longitudinal electric field, subject to the following assumptions:

a. The beam is immersed in a longitudinal dc magnetic field that is intense enough so that electron motion transverse to the beam propagation direction can be neglected. This assumption allows us to consider electron motion along one direction only, which we call the y-direction.

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[1] *F. S. Rusin and G. D. Bogomolov, Soviet Patent 195557 (1965).*

[2] *N. Taguchi and S. Ono, Report of Research Group on Electron Devices, Sendai, Japan (March 1964).*

[3] *V. P. Shestopalov, Diffraction Electronics, Khar'kov, Moscow (1976) (Trans. by U.S. Joint Publications Service, April 1978).*

[4] *S. J. Smith and E. M. Purcell, Phys. Rev., 92 (1953), 1069.*

[5] *R. P. Leavitt, Characteristics of Orotron Oscillation and Amplification, I. Power and Frequency Characteristics, Harry Diamond Laboratories HDL-TR 1899 (July 1979).*

[6] *M. B. Tsevtlin, G. A. Bernashevskiy, V. D. Kotov, and I. T. Tsitson', Radio Eng. Elect. Phys., 22 (1977), 132.*

[7] *K. Mizuno, S. Ono, and Y. Shibata, IEEE Trans. Electron Dev.,* ED-20 *(1973), 749.*

* *D. E. Wortman and R. P. Leavitt, Near Millimeter Wave Orotron Research Study, Harry Diamond Laboratories (draft).*

[5] *R. P. Leavitt, Characteristics of Orotron Oscillation and Amplification, I. Power and Frequency Characteristics, Harry Diamond Laboratories HDL-TR-1899 (July 1979).*

[8] *F. A. Hopf, P. Meystre, G. T. Moore, and M. O. Scully, in Physics of Quantum Electronics, Vol. 5, Novel Sources of Coherent Radiation, Addison-Wesley Publishing Co., Inc., Reading, MA (1978).*

[9] *P. Sprangle and A. Drobot, IEEE Trans. Microwave Theory Tech.,* MTT-25 *(1977), 528.*

b. The electron beam accelerating voltage is low enough so that relativistic effects can be neglected. This assumption allows us to use the nonrelativistic equations of motion to analyze the electron beam propagation.

c. The density of electrons in the beam is low enough so that the forces of repulsion between the electrons can be neglected. This assumption avoids the need to use a self-consistent solution of Maxwell's equations and the dynamical equations.

d. The longitudinal electric field in the resonator is weak enough so that it may be considered as a perturbation on the freely propagating electron beam. This assumption allows us to make a perturbation expansion of the dynamical equations and is consistent with the point of view that regards the orotron as a "free-electron laser of the Smith-Purcell type."[8]

One way to approach the problem of bunching in orotrons is to consider the behavior of a single electron under the influence of the longitudinal rf field of the open resonator.[6,7] This approach is cumbersome and leads to the same results as those given here. In this work, we approach the problem via the distribution function approach, which is equivalent to the previous method. Furthermore, the distribution function approach can deal with increased density systems where the single electron approach would be hopelessly complicated.

The distribution function describes the density of electrons in coordinate and velocity space. To be more precise, if f(y,v,t) is the distribution function, then f(y,v,t) dy dv is the number of electrons having coordinates between y and y + dy and velocities between v and v + dv at time t. In terms of the distribution function, the charge density of the electron beam at position y at time t is

$$\rho(y,t) = -e \int_{-\infty}^{\infty} dv\ f(y,v,t) \quad , \qquad (1)$$

and the current density at position y at time t is

$$j(y,t) = -e \int_{-\infty}^{\infty} dv\ v\ f(y,v,t) \quad . \qquad (2)$$

In equations (1) and (2), e is the magnitude of the electron charge. The dynamical equation of motion for the electron beam in terms of its distribution function is given by the Boltzmann equation:[10]

$$\frac{\partial f}{\partial t} + v\frac{\partial f}{\partial y} - \frac{e}{m} E \frac{\partial f}{\partial v} = 0 \quad , \qquad (3)$$

where m is the electron mass and E = E(y,t) is the longitudinal rf electric field. By integrating equation (3) over velocities and using equations (1) and (2), the reader may verify that the charge conservation equation

$$\frac{\partial j}{\partial y} + \frac{\partial \rho}{\partial t} = 0$$

is satisfied.

In deriving equation (3), we have used assumptions a to c. That is,

a. Equation (3) is one-dimensional.
b. Equation (3) is nonrelativistic.
c. Equation (3) does not contain the electronic repulsion forces.

Within these constraints, equation (3) exactly represents the electron motion.

Assumption d is incorporated into our discussion by expanding equation (3) in a power series in the strength of the electric field. We introduce the strength parameter λ and let $E \rightarrow \lambda E$ in equation (3). Then we expand the distribution function in a

[6] M. B. Tseytlin, G. A. Bernashevskiy, V. D. Kotov, and I. T. Tsitson', Radio Eng. Elect. Phys., 22 (1977), 132.

[7] K. Mizuno, S. Ono, and Y. Shibata, IEEE Trans. Electron Dev., ED-20 (1973), 749.

[8] F. A. Hopf, P. Meystre, G. T. Moore, and M. O. Scully, in Physics of Quantum Electronics, Vol. 5, Novel Sources of Coherent Radiation, Addison-Wesley Publishing Co., Inc., Reading, MA (1978).

[10] R. C. Davidson, Methods in Nonlinear Plasma Theory, Academic Press, New York (1972).

power series in λ:

$$f(y,v,t) = \sum_{n=0}^{\infty} \lambda^n f_n(y,v,t) \quad . \quad (4)$$

This expansion is then substituted into equation (3), and the coefficient of λ^n is equated to zero. The result is, for $n = 0$,

$$\frac{\partial f_0}{\partial t} + v \frac{\partial f_0}{\partial y} = 0 \quad , \quad (5)$$

and, for $n > 0$,

$$\frac{\partial f_n}{\partial t} + v \frac{\partial f_n}{\partial y} = \frac{e}{m} E \frac{\partial f_{n-1}}{\partial v} \quad . \quad (6)$$

We choose the zeroth-order distribution function to be a function of velocity only, to correspond to the physical situation of a uniform dc electron beam. The normalization of f_0 is chosen such that the zeroth-order current density is

$$j_0 = -e \int_{-\infty}^{\infty} dv \; v \; f_0(v) = I_0 / A_0 \quad , (7)$$

where I_0 is the total beam current and A_0 is the beam cross-sectional area. This choice of the zeroth-order distribution function satisfies equation (5) since f_0 is independent of y and t.

The longitudinal rf electric field is chosen to be monochromatic and can be written in the form

$$E(y,t) = E_0(y)e^{-i\omega t} + \text{c.c.} \quad , \quad (8)$$

where ω is the angular frequency of the field and c.c. means complex conjugate. We wish to consider the first-order contribution to the distribution function; equation (6) becomes, for $n = 1$,

$$v \frac{d\alpha_1}{dy} - i\omega\alpha_1 = \frac{e}{m} E_0(y) f_0'(v) \quad , \quad (9)$$

where $f_0'(v) = df_0/dv$. In writing equation (9), we have introduced the first-order distribution amplitude α_1 via

$$f_1(y,v,t) = \alpha_1(y,v)e^{-i\omega t} + \text{c.c.} \quad (10)$$

Equation (9) is a linear, first-order, ordinary differential equation, which may be solved by standard methods.[11] The result is

$$\alpha_1(y,v) = \frac{e}{mv} e^{i\omega y/v} f_0'(v) \times \int_{-\infty}^{y} dy' \; e^{-i\omega y'/v} E_0(y') \quad , \quad (11)$$

where we have used the condition $\alpha_1(-\infty,v) = 0$. Substituting equation (11) into equation (10) gives the first-order distribution function.

The first-order contribution to the current density is obtained from equation (2). The velocity integral in equation (2) may be integrated by parts, giving the result

$$j_1(y) = -\frac{ie^2\omega}{m} \int_{-\infty}^{\infty} \frac{dv}{v^2} e^{i\omega y/v} f_0(v) \times \int_{-\infty}^{y} dy' \; (y - y')e^{-i\omega y'/v} E_0(y')e^{-i\omega t} + \text{c.c.} \quad (12)$$

The power radiated by the electron beam is given by

$$P_{rad} = - \left\langle \int jE \, d\tau \right\rangle \quad , \quad (13)$$

where the integral is over the volume of the electron beam and the brackets indicate a time average. Substituting equations (8) and (12) into equation (13) yields the result

$$P_{rad} = \frac{e^2\omega}{m} \frac{d}{d\omega} \left[\int_{-\infty}^{\infty} \frac{dv}{v} f_0(v) \times \int_{\text{beam}} dx \, dz \int_{-\infty}^{\infty} dy \int_{-\infty}^{y} dy' \times e^{i\omega(y - y')/v} E_0^*(y)E_0(y') + \text{c.c.} \right] \quad , \quad (14)$$

where the integral over the transverse coordinates x and z is over the cross-sectional area of the beam. The longitudinal field, $E_0(y)$, depends implicitly on x and z in addition to y.

[11] H. W. Reddick, *Differential Equations*, John Wiley and Sons, Inc., New York (1943).

Equation (14) allows the computation of the power radiated by the electron beam provided that the distribution of electric field strength, $E_0(y)$, is known. It is more convenient, however, to characterize the electric field in terms of its Fourier transform:

$$\tilde{E}_0(p) = \int_{-\infty}^{\infty} e^{-ipy} E_0(y)\, dy \quad . \quad (15)$$

In terms of the Fourier transform, the double integral in equation (14) becomes

$$\int_{-\infty}^{\infty} dy \int_{-\infty}^{y} dy'\, e^{i\omega(y-y')/v} E_0^*(y) E_0(y') = \frac{1}{2}\left|\tilde{E}_0\left(\frac{\omega}{v}\right)\right|^2 - \frac{i}{2\pi} P \int_{-\infty}^{\infty} \frac{dp\ |\tilde{E}_0(p)|^2}{p - \frac{\omega}{v}} \quad , \quad (16)$$

where P means principal value. By using equation (16), equation (14) becomes

$$P_{rad} = \frac{e^2\omega}{m} \int_{-\infty}^{\infty} \frac{dv}{v} f_0(v) \times \int_{beam} dx\, dz\, \frac{d}{d\omega}\left|\tilde{E}_0\left(\frac{\omega}{v}\right)\right|^2 \quad . \quad (17)$$

In what follows, we are interested in the radiation from an electron beam in which all of the electrons are traveling at a velocity v_0. By using the normalization given by equation (8), the distribution function is

$$f_0(v) = -\frac{I_0}{ev_0A_0}\,\delta(v - v_0) \quad , \quad (18)$$

and the power radiated at frequency ω is

$$P_{rad} = -\frac{e\omega I_0}{mv_0^2A_0} \times \int_{beam} dx\, dz\, \frac{d}{d\omega}\left|\tilde{E}_0\left(\frac{\omega}{v_0}\right)\right|^2 \quad . \quad (19)$$

This is the main result of this section; it allows the computation of the power radiated by an electron beam immersed in a longitudinal monochromatic electric field. In the next section, we consider the orotron, in which the electric field is quasi-periodic.

3. RADIATION BY ELECTRON BEAM NEAR PERIODIC STRUCTURE

In applying equation (19) to radiation in an orotron, we must account for the detailed structure of the electromagnetic field distribution. Consider the schematic diagram of the orotron shown in figure 1. The electron beam passes close to the reflecting diffraction grating, and the interaction of the beam with the longitudinal rf field of the resonator (as modified by the grating) causes the beam to radiate, as shown in section 2.

The electromagnetic field of the open resonator of the orotron may be considered as a natural mode of the resonator, with a fine structure imposed by the grating. This fine structure causes the radiation characteristic of the orotron.

Consider a plane wave of frequency ω incident normally on the surface of the grating, as shown in figure 2. Let the period of the grating be ℓ. The electromagnetic field may be characterized by the x-component of magnetic induction, given above the grating by*

$$H_x = H_0\left[2\cos k(z - z_0) + \sum_{r=1}^{\infty} a_r \times \cos\left(\frac{2r\pi y}{\ell}\right) e^{-\Gamma_r z}\right] e^{-i\omega t} + \text{c.c.} \quad , \quad (20)$$

where $k = \omega/c$, c is the speed of light, H_0 is the amplitude of the incoming wave, and

$$\Gamma_r = \left[\left(\frac{2r\pi}{\ell}\right)^2 - k^2\right]^{1/2} \quad . \quad (21)$$

*D. E. Wortman and R. P. Leavitt, *Near Millimeter Wave Orotron Research Study*, Harry Diamond Laboratories (draft).

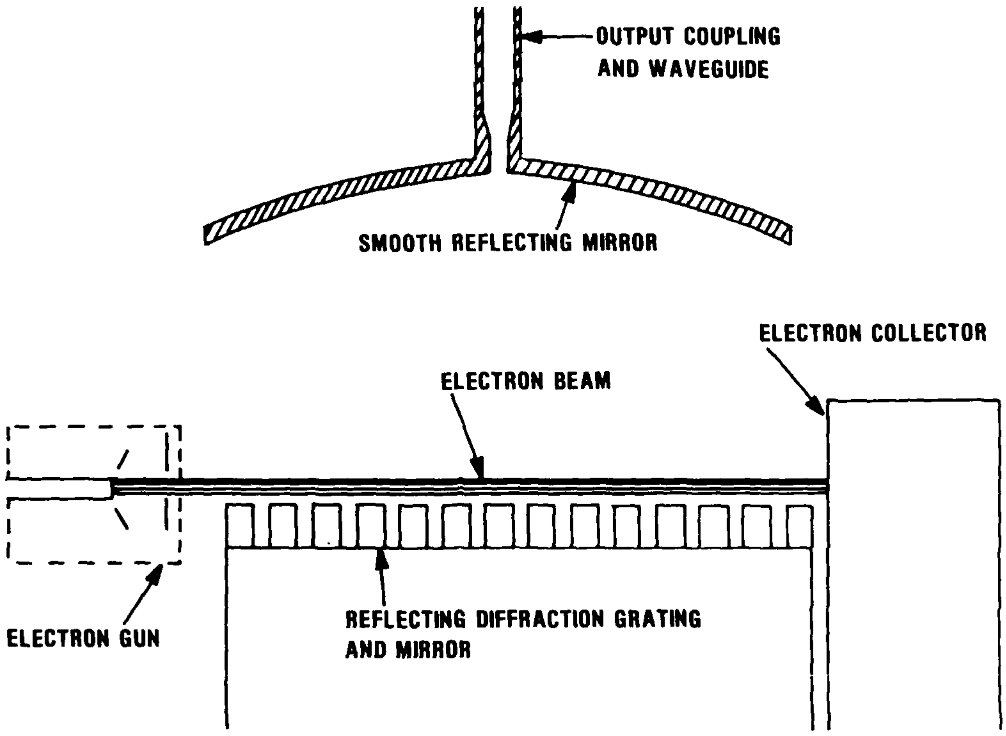

Figure 1. Schematic diagram of orotron.

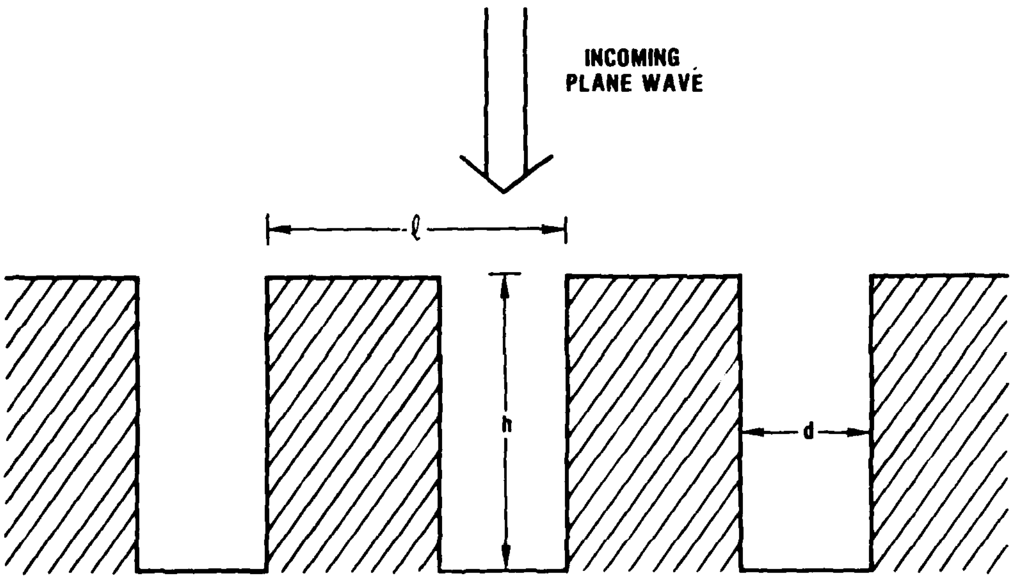

Figure 2. Geometry for calculation of electro-magnetic field distribution of plane wave incident on reflecting diffraction grating.

Equation (20) represents the most general free-space solution of Maxwell's equations corresponding to an incoming plane wave that satisfies the periodicity requirement

$$H_x\,(y + \ell, z) = H_x\,(y, z) \qquad (22)$$

imposed by the periodic grating. The unknown coefficients, a_r in equation (20), are determined by solving the appropriate boundary value problem at the grating surface. We do not concern ourselves with the solution of the boundary value problem here, but we consider the coefficients a_r as being known.* The y-component of the electric field is, from Maxwell's equations,

$$E_y = \frac{i}{\omega\varepsilon_0}\frac{\partial H_x}{\partial z} =$$

$$- \frac{iH_0}{c\varepsilon_0}\left[2 \sin k\,(z - z_0) + \sum_{r=1}^{\infty} a_r \frac{\Gamma_r}{k} \times \cos\left(\frac{2r\pi y}{\ell}\right) e^{-\Gamma_r z}\right] e^{-i\omega t} + \text{c.c.} \quad , (23)$$

where ε_0 is the vacuum permittivity.

The electric field of equation (23) corresponds to an incident plane wave; it must be modified to account for the mode pattern in the open resonator. For a large class of open resonators, the mode pattern assumes the form $f_x(x)f_y(y)$, where f_x and f_y are Gaussians or Gaussians multiplied by Hermite polynomials describing the variation of electromagnetic fields within the resonator. The functions are normalized to unity:

$$\int_{-\infty}^{\infty} dx\, \left|f_x(x)\right|^2 = 1 \quad , (24a)$$

$$\int_{-\infty}^{\infty} dy\, \left|f_y(y)\right|^2 = 1 \quad . (24b)$$

In terms of these, the longitudinal electric field of equation (8) becomes

$$E_0(y) = -\frac{iH_0}{c\varepsilon_0}\left[2 \sin k\,(z - z_0) + \sum_{r=1}^{\infty} a_r \frac{\Gamma_r}{k} \cos\left(\frac{2r\pi y}{\ell}\right) e^{-\Gamma_r z}\right] \times f_x(x)f_y(y) \quad . \qquad (25)$$

This is the electric field that is used in equation (19) to compute the radiated power.

In computing the Fourier transform of this electric field, it is necessary to make a few approximations whose validity depends on the fact that the envelope function $f_y(y)$ is slowly varying compared with the variation of the terms in brackets in equation (25). The function varies slowly if the dimension of the resonator mode along y is very large in comparison with the period of the grating. By using equation (15), the Fourier transform becomes

$$\tilde{E}_0(p) = -\frac{iH_0}{c\varepsilon_0}\sum_{N=-\infty}^{\infty}\int_{-\ell/2}^{\ell/2} dy \times e^{-ip(y + N\ell)} f_x(x)f_y(y + N\ell) \times \left[2 \sin k\,(z - z_0) + \sum_{r=1}^{\infty} a_r \frac{\Gamma_r}{k} \cos\left(\frac{2r\pi y}{\ell}\right) e^{-\Gamma_r z}\right] \quad , (26)$$

where we have broken the y-axis up into segments of length ℓ and used the periodicity condition, equation (22).

It is known that radiation from the grating should peak at frequencies near those given by the Smith-Purcell condition[4,*]

*The determination of the coefficients a_r will be discussed in a forthcoming report in this series

[4] S. J. Smith and E. M. Purcell, Phys. Rev., 92 (1953), 1069

* D. E. Wortman and R. P. Leavitt, Near Millimeter Wave Orotron Research Study, Harry Diamond Laboratories (draft).

$$\omega = \frac{2\pi r v_0}{\ell} \quad . \tag{27}$$

Therefore, referring to equation (19), we shall consider the Fourier transform near $p \sim 2\pi r/\ell$, where r is an integer.* Assuming that f_y does not vary appreciably over an interval of length ℓ, we have

$$\widetilde{E}_0 \left(\frac{2\pi r + \delta}{\ell}\right) =$$

$$- \frac{iH_0}{c\varepsilon_0} f_x(x) \sum_{N=-\infty}^{\infty} f_y(N\ell)\, e^{-i\delta N} \times$$

$$\int_{-\ell/2}^{\ell/2} e^{-i2\pi r y/\ell} \left[2 \sin k (z - z_0) + \right. \tag{28}$$

$$\left. \sum_{r'=1}^{\infty} a_{r'} \frac{\Gamma_{r'}}{k} \cos\left(\frac{2r'\pi y}{\ell}\right) e^{-\Gamma_{r'} z} \right] \quad .$$

The integral selects out the term with r' = r; the remaining sum can be approximated by an integral. The result is

$$\widetilde{E}_0 \left(\frac{2\pi r + \delta}{\ell}\right) =$$

$$- \frac{iH_0 a_r \Gamma_r}{2c\varepsilon_0 k} f_x(x)\, e^{-\Gamma_r z}\, \widetilde{f}_y \left(\frac{\delta}{\ell}\right) \quad , \tag{29}$$

where

$$\widetilde{f}_y(p) = \int_{-\infty}^{\infty} dy\, e^{-ipy} f_y(y) \tag{30}$$

is the Fourier transform of the envelope function, f_y. We find the radiated power to be, from equation (19),

$$P_{rad} = - \frac{e \ell I_0 \left|H_0\right|^2 a_r^2 \Gamma_r^2}{4 m v_0^3 \omega A_0 \varepsilon_0^2}$$

$$\int_{beam} dx\, dz \left|f_x(x)\right|^2 e^{-2\Gamma_r z} \times$$

$$\frac{d}{d\delta} \left| f_y\left(\frac{\delta}{\ell}\right)\right|^2 \quad , \tag{31}$$

at a frequency $\omega = (v_0/\ell)(2\pi r + \delta)$.

Specifically, equation (31) may be evaluated for the case of a TEM_{20} mode in the resonator.* In this case, we have

$$f_x(x) = \left[\frac{1}{4 (2\pi)^{1/2} w_x}\right]^{1/2} \times$$

$$H_2 \left(\frac{2^{1/2}}{w_x}\right) e^{-x^2/w_x^2} \quad , \tag{32}$$

where H_2 is the second-degree Hermite polynomial, and

$$f_y(y) = \left[\frac{2}{(2\pi)^{1/2}}\right]^{1/2} e^{-y^2/w_y^2} \quad . \tag{33}$$

The width parameter, w_x, is chosen equal to the width of the electron beam, D. The integral over the electron beam cross section in equation (31) is then

$$\int_{beam} dx\, dz \left| f_x(x)\right|^2 e^{-2\Gamma_r z} =$$

$$0.2 \frac{e^{-2\Gamma_r a}}{2\Gamma_r} \left(1 - e^{-2\Gamma_r \triangle}\right) \quad , \tag{34}$$

where a is the height of the beam above the grating and $\triangle$ is the beam thickness ($A_0 = D\triangle$). The Fourier transform of $f_y(y)$ can be evaluated in a straightforward manner to give

$$\widetilde{f}_y \left(\frac{\delta}{\ell}\right) = (2\pi)^{1/4} (w_y)^{1/2} e^{-w_y^2 \delta^2/4\ell^2} \quad . \tag{35}$$

The quantity $\left|H_0\right|^2$ in equation (31) may be eliminated in terms of the total energy in the open resonator,

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* We refer to an orotron radiating at frequencies near $2\pi r v_0/\ell$ as an orotron operating on the rth harmonic.

* The TEM_{20} mode is the mode most commonly encountered in practice.

$$U = \frac{1}{2} < \int (\varepsilon_0 E^2 + \mu_0 H^2)\, d\tau > \quad ,(36)$$

where the integral is over the volume of the open resonator and μ_0 is the vacuum permeability. Neglecting the contribution of the evanescent waves in equation (25), we obtain

$$U = \frac{4LH_0^2}{\varepsilon_0 c^2} \quad , \quad (37)$$

where L is the mirror spacing in the open resonator and c is the speed of light *in vacuo.* Inserting equations (34), (35), and (37) into equation (31) gives the result

$$P_{rad} = - \frac{0.2(2\pi)^{1/2} eI_0 U a_r^2 \Gamma_r^2 c^2}{16\ mv_0^3 \omega DL\varepsilon_0} \times$$

$$\frac{e^{-2\Gamma_r a}}{2\Gamma_r \triangle}\left(1 - e^{-2\Gamma_r \triangle}\right) \frac{w_y^3}{\ell}\ \delta\ e^{-w_y^2 \delta^2 / 2\ell^2} \ .(38)$$

The orotron actually radiates at a frequency given by the peak of equation (38); the peak is at

$$\frac{w_y \delta}{\ell} = -1 \quad (39)$$

or

$$\omega = \frac{v_0}{\ell}\left(2\pi r - \frac{\ell}{w_y}\right) \quad . \ (40)$$

Thus, the peak output power from the device is

$$P_{rad}(\max) = 0.0095\ |I_0|\ U \frac{ea^2 \Gamma_r^2 c^2 w_y^2}{mv_0^3 \omega DL\varepsilon_0} \times$$

$$\left[\frac{e^{-2\Gamma_r a}}{\Gamma_r \triangle}\left(1 - e^{-2\Gamma_r \triangle}\right)\right] \quad . \ (41)$$

The starting current for orotrons may be determined by requiring that the power gain given by equation (41) equal the power loss in the resonator, given by

$$P_{loss} = \frac{U\omega}{Q} \quad . \quad (42)$$

Thus, the starting current is given by

$$I_s(\min) = \frac{mv_0^3 \omega^2 DL\varepsilon_0}{0.0095 ea_r^2 \Gamma_r^2 c^2 w_y^2 Q} \times$$

$$\left[\frac{e^{-2\Gamma_r a}}{\Gamma_r \triangle}\left(1 - e^{-2\Gamma_r \triangle}\right)\right]^{-1} \quad . (43)$$

Equation (43) has been evaluated for the orotron being built at HDL to operate in the fundamental mode (r = 1) at 75 gigahertz. The electromagnetic boundary problem in equation (20) has been solved* with ℓ = 0.4 millimeter to give

$$a_1 = 0.2796,$$

$$\Gamma_1 = 1.563 \times 10^4/\text{meter}. \quad (44)$$

We consider the case of a beam that grazes the grating; therefore, a = 0. The other parameters of the system are

$m = 9.1 \times 10^{-31}$ kilogram,

$v_0 = 3 \times 10^7$ meters/second,

$\omega = 2\pi \times 75 \times 10^9$/second,

$a \simeq 0$,

$\triangle = 0.03$ centimeter,

$D = 1$ centimeter,

$L = 2$ centimeters,

$\varepsilon_0 = 8.854 \times 10^{-12}$ farad/meter,

$e = 1.6 \times 10^{-19}$ coulomb,

$c = 3 \times 10^8$ meters/second,

$w_y = 1$ centimeter,

$Q = 5000$ (estimated). (45)

**The determination of the coefficients a_r will be discussed in a forthcoming report in this series.*

Using these parameters, we obtain

$$I_s = 34.5 \text{ milliamperes.} \tag{46}$$

4. ELECTRONIC TUNING CHARACTERISTICS OF OROTRON

In a previous report of this series,[5] we showed that the frequency characteristics of an orotron could be derived from an equivalent circuit representation. The electron beam was represented in this circuit by a negative conductance that, if sufficiently large in magnitude, acted as a feedback element leading to oscillation of the circuit. This model predicted the correct behavior of the orotron with regard to the threshold condition and the behavior of the output line width below and above threshold. One of the parameters in the model, the threshold current, I_s, has been calculated in this report (sect. 3) from first principles. However, the electron beam is not actually represented by a frequency-independent conductance, as was assumed in the previous report.

Consider equation (14) for the radiated power. The real part of equation (16) represents the effect of that part of the current that is in phase with the applied electric field; this is the only part that contributes to the radiated power. However, the imaginary part (which does not contribute to the time-averaged power) representing the out-of-phase (by 90 deg) current component can be considered as being due to an "imaginary conductance" (rigorously, a susceptance). Therefore, the overall effect of the electron beam is to introduce a frequency-dependent complex admittance (rather than a simple conductance, as was assumed earlier)[5] into the equivalent circuit.

We consider an admittance that is proportional to the derivative of equation (16):

$$Y_{beam} = \beta \frac{d}{d\omega}\left[\left|\widetilde{E}_0\left(\frac{\omega}{v_0}\right)\right|^2 - \frac{1}{\pi} P \int_{-\infty}^{\infty} \frac{dp\,|\widetilde{E}_0(p)|^2}{p - \left(\frac{\omega}{v_0}\right)}\right] , \tag{47}$$

where β is a proportionality constant and we have considered the sharply peaked distribution function given by equation (18). We have chosen the sign in equation (47) so that real (Re) $Y_{beam} > 0$ corresponds to radiation.

Section 3 is concerned with the evaluation of the first term in equation (47). Since at this point we are interested only in the frequency dependence, we can represent the result of that calculation by a "normalized conductance":

$$\mathrm{Re}\ y_{beam} = -\frac{I_0}{I_s(\min)}\, u\, e^{(1-u^2)/2} \ , \tag{48}$$

where

$$u = \frac{w_y \delta}{\ell} = w_y\left(\frac{\omega}{v_0} - \frac{2r\pi}{\ell}\right) .$$

This result is obtained by squaring equation (35) and differentiating with respect to δ. The normalization is chosen so that, at $u = -1$ (where Re y_{beam} attains its maximum) and at a current $I_0 = I_s$ (min), the beam conductance Re $y_{beam} = 1$.

From equation (47), the imaginary (Im) part of the beam admittance is given by

$$\mathrm{Im}\ y_{beam}(u) = -\frac{1}{\pi}\int_{-\infty}^{\infty} \frac{du'\ \mathrm{Re}\ y_{beam}(u')}{u' - u} \ . \tag{49}$$

This integral may be performed easily by using equation (48); the result is

$$\mathrm{Im}\ y_{beam}(u) = \left(\frac{2e}{\pi}\right)^{1/2} \frac{I_0}{I_s(\min)}\, F\left(1;\frac{1}{2};-\frac{u^2}{2}\right) \ . \tag{50}$$

[5] R. P. Leavitt, *Characteristics of Orotron Oscillation and Amplification, I Power and Frequency Characteristics*, Harry Diamond Laboratories HDL-TR 1899 (July 1979).

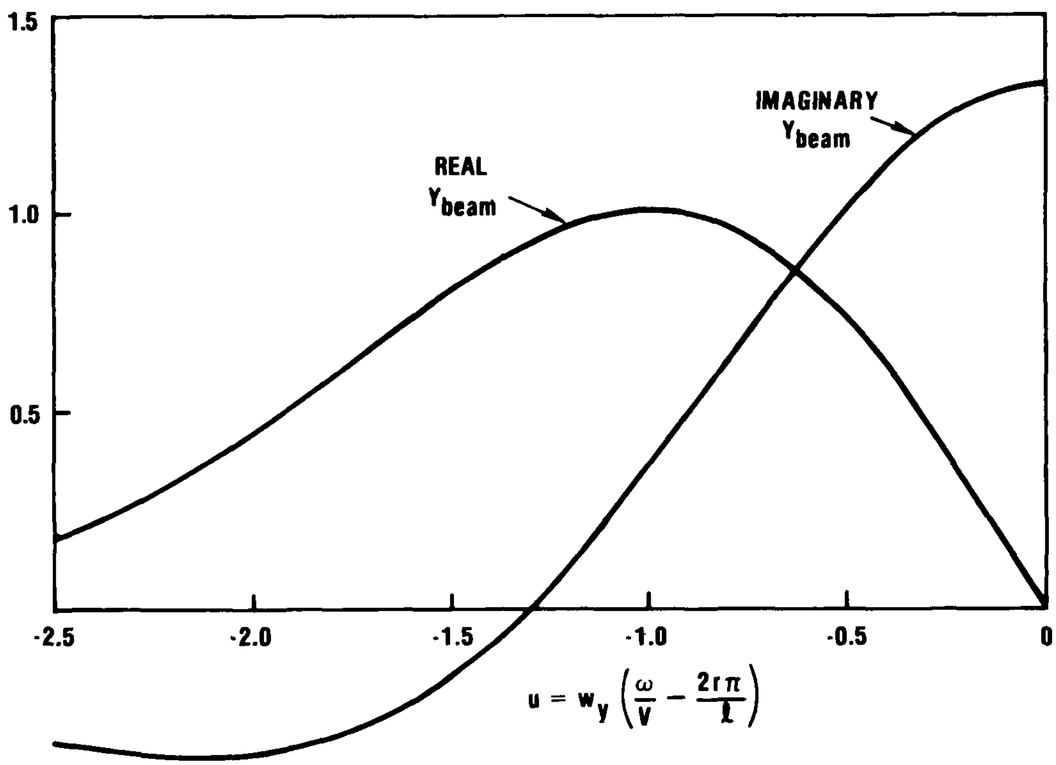

Figure 3. Real and imaginary parts of normalized electron beam admittance at $I_0 = I_s$ (min).

where F is a confluent hypergeometric function. The real and imaginary parts of y_{beam} are shown in figure 3 in the region $u < 0$, where the beam radiates.

Consider the orotron equivalent circuit shown in figure 4, which is the circuit given earlier[5] generalized in accordance with the above in terms of normalized admittances. The overall admittance of the circuit (normalized) is given by

$$y = y_0 + g_{out} - y_{beam} \quad , \qquad (51)$$

where y_0 is the normalized admittance of the open resonator and g_{out} is the normalized conductance representing the output coupling. By using the results of reference 5, equation (51) may be written as

$$y = \left[1 - \mathrm{Re}\ y_{beam}(u)\right] - i\left[2Q\,\frac{\omega - \omega_0}{\omega_0} + \mathrm{Im}\ y_{beam}(u)\right] \quad , (52)$$

where Q is the quality factor of the passive open resonator and ω_0 is the resonant frequency of the resonator. The normalization is chosen so that $\mathrm{Re}\ y_0 + g_{out} = 1$.

If we consider the orotron as an ideal oscillator (that is, we neglect the noise source in fig. 4), the condition for oscillation is that both the real and the imaginary parts of equation (52) must vanish. That is,

$$\mathrm{Re}\ y_{beam} = 1 \qquad (53)$$

for the real part and

$$\omega = \omega_0 \left[1 - \frac{\mathrm{Im}\ y_{beam}(u)}{2Q}\right] \qquad (54)$$

[5] *R. P. Leavitt, Characteristics of Orotron Oscillation and Amplification, I. Power and Frequency Characteristics, Harry Diamond Laboratories HDL-TR-1899 (July 1979).*

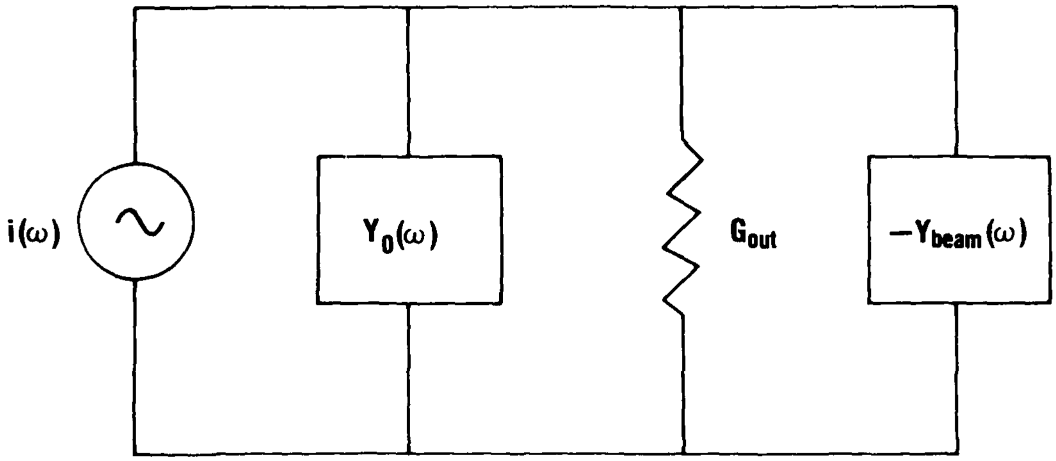

Figure 4. Equivalent circuit of orotron.

for the imaginary part. By using equation (48), equation (53) implies a starting current of

$$I_s = -\frac{I_s(\text{min})}{u} e^{(u^2-1)/2} \quad , u < 0 \quad . (55)$$

The starting current is shown as a function of the frequency variable u in figure 5.

Since

$$u = w_y\left(\frac{\omega}{v_0} - \frac{2r\pi}{\ell}\right) ,$$

equation (54) must be solved self-consistently for the frequency of orotron oscillation. However, if the Q of the open resonator is sufficiently large, we may replace u by

$$u_0 = w_y\left(\frac{\omega_0}{v_0} - \frac{2r\pi}{\ell}\right)$$

on the right-hand side of equation (54). Using equation (55) in equation (50) and substituting the result into equation (54), we obtain

$$\omega = \omega_0 \times$$

$$\left[1 + \frac{e^{u_0^2/2}}{(2\pi)^{1/2} Q u_0} F\left(1;\frac{1}{2};-\frac{u_0^2}{2}\right)\right] \quad . (56)$$

The second term in equation (56) is a very small correction to the first, since $Q \sim 5000$ and $u_0 \sim -1$. Therefore, the output frequency is very close to the resonant frequency of the passive open resonator. However, this correction determines the electronic tuning characteristics of the orotron as follows. The change in output frequency with a change in accelerating voltage V is

$$\frac{d\omega}{dV} = \frac{\omega_0}{(2\pi)^{1/2} Q} \frac{dv_0}{dV} \frac{du_0}{dv_0} \frac{d}{du_0} \times$$

$$\left[\frac{e^{u_0^2/2}}{u_0} F\left(1;\frac{1}{2};-\frac{u_0^2}{2}\right)\right] \quad . (57)$$

Using

$$V = \frac{1}{2}\frac{m}{e} v_0^2 \quad ,$$

$$\frac{dV}{dv_0} = \frac{2}{v_0} V \quad ,$$

and

$$\frac{du_0}{dv_0} = -\frac{w_y \omega_0}{v_0^2} \quad .$$

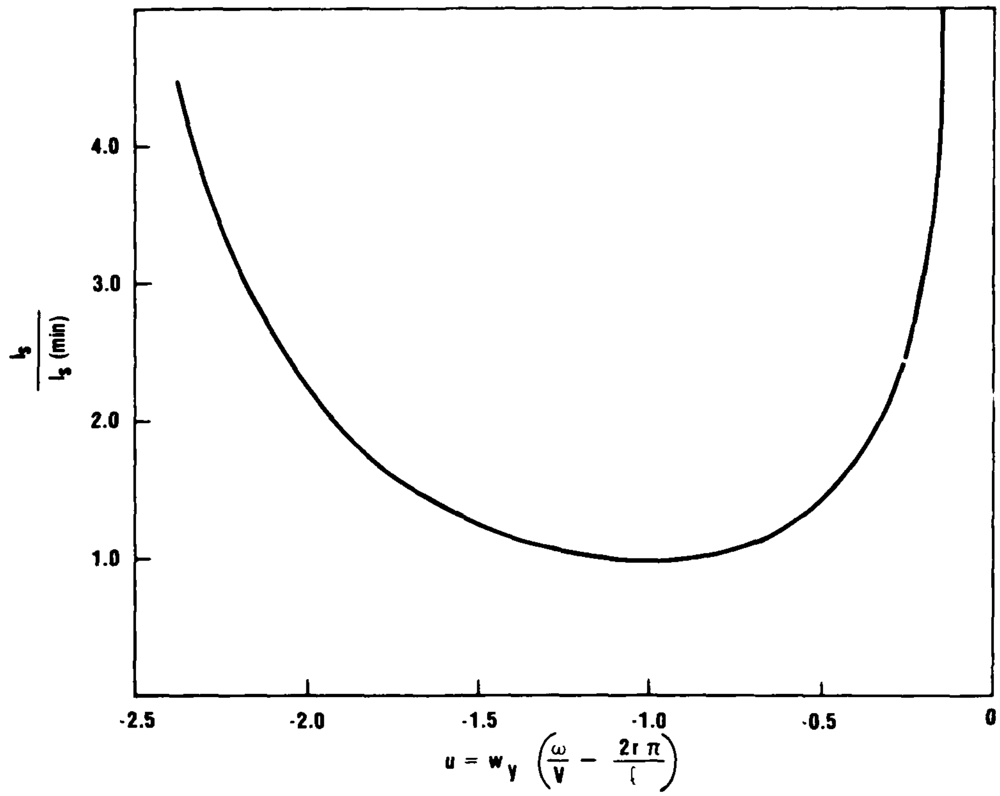

Figure 5. Variation of starting current in orotron.

we obtain (using the properties of the confluent hypergeometric function)[12]

$$\frac{d\omega}{dV} = \frac{w_y \omega_0^2}{(8\pi)^{1/2}\, Q V v_0 u_0^2}\, F\left(\frac{1}{2};\frac{1}{2};\frac{u_0^2}{2}\right) . \quad (58)$$

We are primarily interested in the evaluation of equation (58) in the region near $u_0 = -1$, where the starting current of equation (55) is a minimum. The hypergeometric function in equation (58) has the value $e^{1/2}$ at this point; and so equation (58) becomes

$$\left(\frac{d\omega}{dV}\right)_0 = \left(\frac{e}{8\pi}\right)^{1/2} \frac{w_y \omega_0^2}{Q V v_0} . \quad (59)$$

We consider the evaluation of equation (59) for the HDL 75-gigahertz orotron experiment. Using

$$\omega_0 = 2\pi \times 75 \times 10^9/\text{second},$$

$$w_y = 1 \text{ centimeter},$$

$$V = 2500 \text{ volts},$$

$$v_0 = 3 \times 10^7 \text{ meters/second},$$

$$Q = 5000 \text{ (estimated)}, \quad (60)$$

we obtain for the electronic tuning transconductance

$$\frac{df}{dV} = \frac{1}{2\pi} \left(\frac{d\omega}{dV}\right)_0 =$$

$$0.31 \text{ megahertz/volt} \quad . \quad (61)$$

[12] *I. S. Gradshteyn and I. M. Ryzhik, Table of Integrals, Series, and Products, Academic Press, New York (1965).*

5. CONCLUSIONS

We have derived a linear theory of electron bunching in orotrons based on the time evolution of the distribution function representing the electron beam. This approach to bunching has the advantage over conventional approaches[6,7] that it is easily generalizable to a nonlinear theory in which saturation effects are included. The result for the power radiated, which is derived in section 2, applies not only to the orotron, but to any device in which a longitudinal electric field acts as the bunching field. Sections 3 and 4 consider the evaluation of the starting current, I_s, and the electronic tuning transconductance, df/dV. These quantities were evaluated for the planned orotron experiment at 75 gigahertz; the results were I_s = 34.5 milliamperes and df/dV = 0.31 megahertz/volt. It is remarkable that these results agree well with published experimental data taken in the Soviet Union. (For example, in one experimental device similar to the one under construction at HDL, values of I_s = 28 milliamperes and df/dV = 0.37 megahertz/volt were measured.[3])

An extension of the theory to the nonlinear regime is necessary to calculate the actual output power of the orotron as a function of the electron beam current. The formalism introduced in this report will allow such a calculation to be performed in a straightforward manner. This nonlinear theory will also allow us to consider such effects as second harmonic generation in the orotron. This calculation will be the subject of a future report in this series.

In section 3, we consider the electromagnetic field distribution near the reflecting diffraction grating of the orotron. In particular, the unknown coefficients, a_r, in the expansion of the electromagnetic field are determined by solving a boundary value problem near the grating. We have performed these calculations and will report on the results. The grating calculation allows not only the determination of the a_r, but also the calculation of the quality factor Q of the open resonator from first principles.

---

[3] *V. P. Shestopalov, Diffraction Electronics, Khar' kov, Moscow (1976) (Trans. by U.S. Joint Publications Service, April 1978).*

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